Neutrino Emission and Mass Ejection in Quark Novae

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ABSTRACT

We explore the role of neutrinos in a Quark Nova explosion. We study production of neutrinos during this event, their propagation and their interactions with the surrounding quark matter and the neutron-rich envelope. We address relevant physical issues such as the timescale for the initial core collapse, the total energy emitted in neutrinos and their effect on the low density matter surrounding the core. We find that it is feasible that the neutrino burst can lead to significant mass ejection of the nuclear envelope.

Subject headings: quark star – quark nova

1. Introduction

At high baryon density and vanishing pressure, the ground state of bulk matter may not be the most stable isotope of iron (Fe^{56}), but deconfined strange quark matter (SQM)

made up of up, down and strange (u, d, s) quarks (Itoh 1970; Bodmer 1971; Witten 1984; Farhi & Jaffe 1984). In that case, once the density for a transition to the (u, d, s) phase is reached in the core of a neutron star (NS), the entire star is contaminated and converted into a *(u,d,s)* star (Haensel et al. 1986; Alcock et al. 1986; Olinto 1987; Olesen & Madsen 1991; Heiselberg, Baym, & Pethick 1991; Glendenning 1992; Cheng & Dai 1996; Anand et al. 1997; Dey et al. 1998; Bombaci & Datta 2000). The conversion could also happen via clustering of Λ-baryons, causing formation of clumps of strange matter, or by seeding and neutrino sparking (Alcock, Farhi & Olinto (1986)), although details of such processes are still debated.

Here we focus on models where the NS core is already converted into (u, d)-quark matter (Alcock et al 1986). Deconfinement might occur during or after supernova explosion provided the central density of the proto-neutron star is high enough to induce phase conversion (e.g., Dai, Peng, & Lu 1995; Xu, Zhang, & Qiao 2000). Figure 1 is a schematic representation of the resulting hybrid star (HS). It is essential that there be a first-order phase transition between neutron matter and (u, d)-quark matter at some critical pressure when the latter becomes favored. Without the phase transition, there will be no conversion.

In the hypothetical scenario which we call Quark-Nova (QN) (Ouyed, Dey & Dey 2002; ODD) the *(u,d)* core of the HS shrinks to the corresponding stable, more compact *(u,d,s)* configuration faster than the overlaying material (the neutron-rich hadronic envelope; Figure 2) can respond, leading to a "gap-like" region of much lower density in between. In this paper, we investigate if it is possible for the neutrinos created in the conversion[1] of neutron matter to SQM to power the ejection of part of the neutron-rich overlaying envelope and the suspended nuclear matter crust of the quark core. Among the issues to consider: (i) Are the neutrinos trapped, and if so, for how long? Since it is the conversion from *(u,d)* to *(u,d,s)* matter that produces neutrinos, the answer depends on the neutrino mean free path just outside the transition surface. (ii) How much energy is deposited by neutrinos in the surrounding unconverted matter? This depends on several factors such as the average neutrino energy, the local temperature and density, and Y_e, the electron fraction of the untransformed neutron star, all of which influence the neutrino mean free path in normal matter. Y_e and the temperature depend in turn on the degree of deleptonization, or the time between the formation of the neutron star and the transition to *(u,d)* matter in the core. If the transition occurs in the protoneutron star (PNS), $T \sim 50$ MeV and $Y_e \sim 0.3$. If the

[1]Neutrino bursts from these conversions have already been investigated in the literature in the context of cold neutron stars and in supernova cores (e.g., Anand et al. 1997 and references therein). In the QN scenario the shrinking *(u,d,s)* core is small enough for neutrinos to escape. As such, our work is an extension of previous studies to include the effects of escaping neutrinos on the overlaying/infalling neutron-rich matter.

transition occurs well after the birth of the NS, $T \sim 50$ keV, $Y_e \sim 0.01$, and neutrinos free stream through the envelope. (iii) How are the neutrino rates affected if *(u,d,s)* matter is in a color superconducting state? This is relevant for the PNS since critical temperatures can be as large as 50 MeV and neutrino emission and absorption is known to be strongly modified in such phases.

Phase transitions into strange matter may cause mass ejection due to a core bounce (Fryer & Woosley, 1998; see also Takahara & Sato, 1988 and Gentile et al. 1993). In this kind of model, a large strange matter core (with a radius of 4-6 km) is formed and the neutrinos are trapped long enough that they cannot efficiently transport energy to the outer layers; in other words, the resulting neutrino wind is not luminous enough to lead to mass ejection. However, according to hydrodynamical simulations the core bounce may cause baryon rich mass ejection with relativistic Γ-factors of the order of $\Gamma \sim 40$, which is much lower than needed for a gamma ray burst. In this work we study the other extreme and neglect the core bounce. We show that in a hybrid star with a small initial core (a radius of 1-2 km) of (u,d) quarks a phase transition into strange matter can lead to a neutrino driven wind that can expel some of the matter of the outer hadronic envelope of the star. A very small core does not produce enough neutrinos to cause mass ejection, whereas in case of a larger core neutrinos are trapped long enough that the wind remains too weak to expel anything. It would be interesting to study both core bounce and the role of neutrinos simultaneously, but this is left as an avenue for future work.

This paper is organized as follows: In Sect. 2 we describe some important features of the QN. In Sect. 3 we identify the dominant neutrino emission and absorption processes and estimate relevant timescales for neutrino diffusion. The full neutrino luminosity and the corresponding mass ejection is calculated in Sect. 4. We present our conclusions and the scope of future investigations in Sect. 5.

2. Quark Nova

The basic idea in the picture we call QN is that the strange matter core, once formed, will be isolated from the rest of the star. Charge neutrality at the *(u,d,s)* surface requires the presence of electrons which are bound to it by strong electrostatic fields that are determined self-consistently in a Thomas-Fermi approximation (Glendenning et al. 1995). This electron gas extends well outside the sharp surface of the *(u,d,s)* object. The resulting potential difference is strong enough to hold up protons against gravity, resulting in a Coulomb gap of 200 fm (Alcock et al., 1986). Neutral particles such as the neutrons constituting most of the envelope/hadronic material will traverse this gap, but protons and other positive ions will

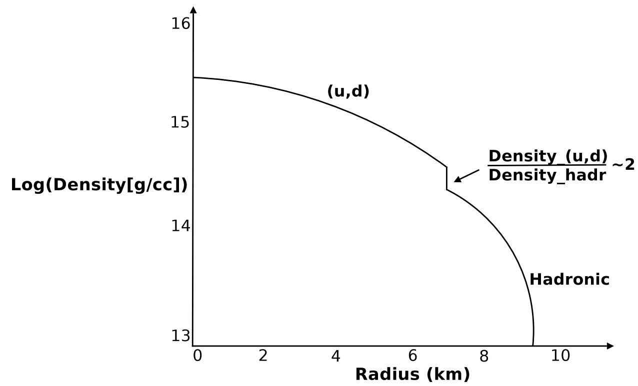

Fig. 1.— **Schematic illustration of Density vs Radius in a hybrid star**: A discontinuity in density, corresponding to the difference between the hadronic and the *(u,d)* core occurs at the radial coordinate where the pressure is equal to that of the mixed phase. (Glendenning 1992; Heiselberg et al. 1993. Also, see Chapter 9 in Glendenning (1997) for more details.

be repelled by the enormous electric field. Following conversion to *(u,d,s)* matter, a large flux of neutrinos is emitted which can deposit energy in the surrounding matter, leading to a mass outflow. Below, we present a first attempt to quantify this scenario.

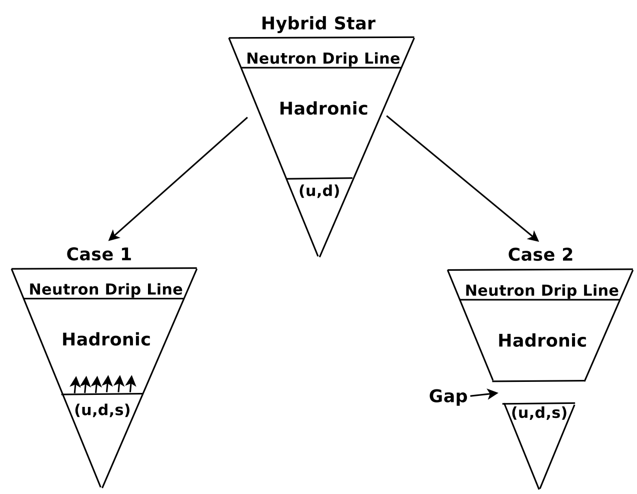

Fig. 2.— **Conversion of a hybrid star into a** *(u,d,s)* **star**: Illustrated are the two possible cases. In the scenario to the right (the QN picture), the *(u,d)* core will shrink faster than the hadronic envelope has time to adjust. In the scenario to the left, the entire HS is converted into *(u,d,s)* without separation of the core and the hadronic envelope.

2.1. Dynamical timescales

The phase transition *(u,d)* $\rightarrow$ *(u,d,s)* inside the HS propagates as a detonation roughly at the speed of sound $c_s = c/\sqrt{3}$. A core few kilometers in radius transforms into a *(u,d,s)* object within $\sim$ 0.1 ms (Lugones et al., 1994). Although weak processes can generate strange quark matter much faster than this, the likelihood that it occurs simultaneously inside the whole object is vanishing.

Interestingly, the hydrodynamical timescale $t_{\rm coll}$ for the shrinking of the core is within the same order of magnitude as the free fall time scale, which is of the form (e.g. Chapter 18.5 in Shapiro & Teukolsky 1983)

$$t_{\rm coll} = \sqrt{\frac{1}{G\rho}} \simeq 0.1\,{\rm ms}\ \sqrt{\frac{10^{15}\,{\rm g\,cm^{-3}}}{\rho_{\rm uds}}}\ , \tag{1}$$

which suggests that the conversion and shrinking of the core occur almost simultaneously.

Given the jump in the density between the HS core and the hadronic envelope (≥ 2), we expect the *(u,d,s)* core to shrink faster than the envelope can respond. Roughly, $t^{uds}_{\rm coll}/t^{env}_{\rm coll} \simeq \sqrt{\rho_{\rm env}/\rho_{\rm uds}} \simeq \sqrt{0.5}$. The spatial gap (see Figure 2), given as $\delta R/R_{\rm ud} = 1 - (\rho_{\rm ud}/\rho_{\rm uds})^{1/3}$, is of the order of a hundred meters and is much larger than the Coulomb gap, implying a large density discontinuity between hadronic and quark matter.

2.2. Energetics

Energy is released in both conversion of baryonic matter into strange matter in the form of latent heat and in the form of gravitational potential energy. The energy per baryon at zero pressure is speculated to be 50 MeV lower in strange matter than in normal matter, and this energy can be released in the phase transition to strange matter (see e.g. Glendenning 1997, p. 351). Integrated over the entire neutron star mass it would be roughly 10^{53} ergs. Moreover, the star shrinks and the gravitational potential energy released is of the same order. As with core-collapse supernova, it is suggested that this energy can be released in the form of neutrinos and/or in the form of luminous ejecta (ODD) depending on the fate of the neutrinos once generated. Energy can also be released in the form of gravitational waves (GWs), where the period of one pulsation is $f = 2\pi\sqrt{R^3/GM} \simeq 4\times 10^{-4}$ s (Cheng & Dai 1998). It is interesting to notice that the pulsation timescale defines the collapse time scale (or $1/(4f) \simeq 1$ ms) if the conversion front into strange matter propagates at least as fast as the star shrinks. In this case, where the deconfinement transition occurs in a dynamical timescale, the GW emission from the phase transition has been examined by

Marranghello, Vasconcellos, & Pacheco (2002). They found that even in the most favorable case corresponding to rapidly rotating (ms period)[2] stars, only 1% of the available energy will be emitted as GWs. These phase transitions excite mainly the radial modes of the star (Sotani, Tominaga, & Maeda, 2002) which can only emit GWs when coupled with rotation (Chau, 1967). As for non-radial modes, and extrapolating from studies done for the case of newly born neutron stars (e.g., Ferrari, Miniutti, & Pons, 2002), these would have comparable frequencies but damping timescales at least one order of magnitude higher than the radial modes coupled to rotation; they would carry away a negligible amount of the energy released. To conclude, the bounds on the magnitude of GWs and on the amount of energy that they can take away during a QN suffer from uncertainties (e.g, values of the viscosity of the *(u,d,s)* matter). How these would exactly affect the dynamics of the QN event is still uncertain at this stage. We now proceed to take a closer look at the neutrinos.

3. Neutrino processes

3.1. Conversion in the core

Once the *(u,d)* density is reached in the core of the HS (following spin-down or accretion), conversion into *(u,d,s)* matter leads to neutrino emission via the following reactions (Iwamoto 1980; Dai et al. 1995):

$$
\begin{aligned}
u + e^- &\leftrightarrow d + \nu_e\,,\\
u + e^- &\leftrightarrow s + \nu_e\,,\\
u + d &\leftrightarrow u + s\,,\\
d &\rightarrow u + e^- + \bar{\nu}_e\,,\\
s &\rightarrow u + e^- + \bar{\nu}_e\,.
\end{aligned}
\tag{2}
$$

In case neutrinos are nondegenerate (low temperatures, when neutrinos can escape), the rates for the first two reactions, which are identical to the rates of the fourth and fifth reactions respectively, were derived by Iwamoto (1982) and Duncan et al. (1983). The first reaction can proceed at an appreciable rate only if strong interactions are also taken into account, in which case it dominates the neutrino emissivity

[2]Recent studies have shown that young *(u,d,s)* stars are stable to the viscosity-driven r-mode instability (Madsen 2000; Andersson, Jones, & Kokkotas 2002). These investigations concluded that the instability cannot develop in *(u,d,s)* stars in any astrophysically relevant temperature window (Gondek-Rosińska, Gourgoulhon, & Haensel 2003).

$$\epsilon_{q\beta} = 1.9 \times 10^{25} \left(\frac{n_b}{n_0}\right) T_9^6 \text{ erg cm}^{-3}\text{s}^{-1}\,, \tag{3}$$

where we have set the strong coupling constant $\alpha_s = 0.1$ and $Y_e = 0.01$, typical of cold quark matter (Iwamoto (1982)), with approximately equal number of u, d, s quarks. T_9 is the temperature in units of 10^9 K and n_B/n_0 is the ratio of baryon density to nuclear matter saturation density $n_0 = 0.16$ fm^{-3}. The rate for the second reaction in eq.(2) is Cabbibo suppressed, hence an order of magnitude slower, and the third reaction is suppressed at tree level in the Standard Model, but can become important at high temperatures (see below).

On the other hand, for temperatures of tens of MeV, at typical (u, d, s) densities, the neutrinos are trapped and degenerate. Emission rates for degenerate neutrinos were derived by Dai et al. (1995), which implies an emissivity

$$\epsilon_{q\beta} = 2 \times 10^{40} \left(\frac{n_b}{n_0}\right)^{2/3} T_{11}^3 \text{ erg cm}^{-3}\text{s}^{-1}\,, \tag{4}$$

where we have chosen $Y_e = 0.3, Y_\nu = 0.1$ (Prakash et al (1995)), $\alpha_s = 0.1$ and T_{11} is the temperature in units of 10^{11} K or ~ 10 MeV. Note that the temperature dependence is only T^3. In comparison, for the third reaction in eq.(2), the numerical factor in front is 4 orders of magnitude smaller. However, its rate is proportional to T^5. Increasing the temperature by a factor of 2 increases the rate by a factor of 2^5, hence this process can become important only at very high temperatures.

The emissivity estimated above does not reflect the true luminosity of the star since neutrinos are trapped as we show in the following section. But first let us estimate first order General Relativity effects on the neutrino emissivities. This can be introduced through redshift factors which can be expressed in terms of the core radius R_1 since in our case the core mass can simply be written as $M_1 = 4\pi/3\rho_{\rm uds}R_1^3$. In Figure 3 we plot the ratio between the redshifted luminosity (effectively seen by the envelope material) to the non-redshifted luminosity as a function of the core size and for two different *(u,d,s)* density. It is clear that a more compact core (higher (u,d,s) density implying higher masses for a given radius) will induce higher redshift. In general cores smaller than 2 km induce almost no redshift. Note that neutrinos, if not trapped, will be further redshifted as they stream outwards since they are subject to an increasing core size. Nevertheless we expect at most the luminosity to be reduced by 50% in our first estimate. While the gravitational redshift will degrade the total energy deposition above the core other effects related to bending of neutrino trajectories might compensate the loss (e.g., Cardall & Fuller 1997).

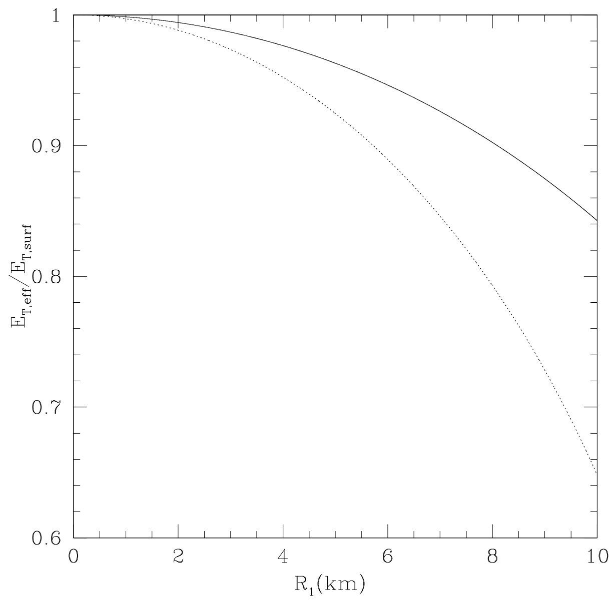

Fig. 3.— The ratio of the redshifted to the non-redshifted emissivity versus core radius (reflecting the compactness parameter). The lower dotted curve is for $\rho_{\rm uds} = 10^{15}$ g cm^{-3} while the upper solid curve is for $\rho_{\rm uds} = 4 \times 10^{14}$ g cm^{-3}. In general, for cores smaller than 2 km general relativistic effects are negligible.

3.2. Diffusion timescales

In this section, we present neutrino opacities that support the claim that neutrinos are essentially thermalized and degenerate in quark matter at MeV temperatures. Besides the inverse processes described in eq.(2), neutrinos can be absorbed by pair annihilation,

$$\nu\bar{\nu} \rightarrow e^+e^- \,. \tag{5}$$

Even at $T \simeq 10$ MeV, where positrons can be created copiously, it is found that the cross-section for this process is 6 orders of magnitude smaller than neutrino scattering off electrons, hence this process can be ignored for estimating the opacity. Absorption by quarks dominates over neutrino-quark scattering at temperatures of tens of MeV and supra-nuclear densities (Steiner et al., 2001).

The main scattering processes are

$$\begin{aligned} e\nu &\rightarrow e\nu \,, \\ q\nu &\rightarrow q\nu \,. \end{aligned} \tag{6}$$

and the neutrino opacity is controlled by scattering against degenerate electrons, for which the cross section is given by (Burrows&Thompson, 2002)

$$\langle\sigma\rangle \simeq 2.5\times10^{-42}\,\mathrm{cm}^2\left(\frac{E_\nu(T_\nu+\mu_e/4)}{(10\,\mathrm{MeV})^2}\right)\ . \tag{7}$$

The mean free path is then

$$\lambda = \frac{1}{n\langle\sigma\rangle} = 500\,\mathrm{cm}\left(\frac{(10\,\mathrm{MeV})^2}{E_\nu(T_\nu+\mu_e/4)}\right)\left(\frac{10^{15}\,\mathrm{g\,cm^{-3}}}{\rho}\right), \tag{8}$$

The corresponding diffusion time we estimate is

$$\begin{aligned}\tau &= (\Delta R)^2/(\lambda c) \simeq 0.1\,\mathrm{s} \\ &\times \left(\frac{\Delta R}{10\,\mathrm{km}}\right)^2\left(\frac{\rho}{10^{15}\,\mathrm{g\,cm^{-3}}}\right)\left(\frac{E_\nu(T_\nu+\mu_e/4)}{(10\,\mathrm{MeV})^2}\right)\ .\end{aligned} \tag{9}$$

Substituting typical numbers $E_\nu \approx T_\nu \simeq 10$ MeV, $\mu_e \approx 100$ MeV and $\rho = \rho_{u,d,s} \sim 10^{15}\,\mathrm{g\,cm^{-3}}$, neutrinos are certainly trapped long enough ($\sim 10-100$ ms) to thermalize and to acquire a black body distribution defined by the *(u,d,s)* effective temperature. Even the neutrinos produced in the outermost layers of 1 km (making up $\sim 27\%$ of the total neutrinos) are diffused out in the timescale of the order of 1 ms. This implies that unless the size of the *(u,d,s)* core is small ($\Delta R << 10$ km), neutrinos are trapped long enough to allow conversion of most of the hadronic envelope material into *(u,d,s)*.

3.3. Conversion of the envelope material

The infalling neutrons traversing the potential barrier will be converted into strange quark matter via the reactions in eq.(2). As mentioned previously, this can happen before diffusing neutrinos reach the contamination front, unless the core radius is small and neutrinos have an energy scale less than about 10 MeV. Note that since the neutron free fall time is much smaller than its free decay lifetime, most neutrons reach the contaminating front of (u, d, s) before decaying to produce neutrinos. Thus, neutrino emission in this time interval is negligible. They will however be produced in flavor equilibrating reactions in strange quark matter. Since the timescale of weak interactions is three orders of magnitude faster than the timescale of neutron free-fall (or, the rate of shrinking of the star), the weak interaction reactions have enough time to equilibrate flavors as well as produce neutrinos (which can also thermalize due to their short mean free path).

4. Neutrino luminosity and mass ejection

For the case where most of the neutrinos are trapped (and thermalized), in the first approximation, the neutrino luminosity can be written as a black-body with the appropriate counting for fermion statistics, leading to

$$\begin{aligned} L_\nu = \frac{7}{8} N_\nu 4\pi R^2 \sigma T_\nu^4 &\simeq 3.2 \times 10^{53}\,\mathrm{erg~s^{-1}} \\ &\times \left(\frac{R}{10~\mathrm{km}}\right)^2 \left(\frac{kT_{\rm eff}}{10~\mathrm{MeV}}\right)^4 , \end{aligned} \tag{10}$$

where $T_{\rm eff}$ is the core temperature and $N_\nu = 3$ is the number of neutrino species. This can be compared to the Eddington luminosity

$$\begin{aligned} L_{\mathrm{Edd},\nu} &\approx 3.0 \times 10^{53}\,\mathrm{erg~s^{-1}} \\ &\times \left(\frac{R}{10~\mathrm{km}}\right)^3 \left(\frac{\rho_{\rm uds}}{10^{15}~\mathrm{gcm^{-3}}}\right) \left(\frac{10\,\mathrm{MeV}}{kT_\nu}\right)^2 . \end{aligned} \tag{11}$$

The above expressions show that $L_\nu > L_{\rm Edd}$ for $T_\nu > T_{\nu,c} = 10$ MeV where the subscript c stands for critical.

The energy release in the phase transition is the same all over the core, so we can safely assume that the temperature is uniform and the average energy of neutrinos in thermal equilibrium would be $T_{\nu,equi.} \simeq 100/4$ MeV (since $n_u \simeq n_d = n_s \simeq n_\nu$). Since $T_{\nu,equi.} > T_{\nu,c.}$, super-Eddington luminosities are achieved. Nevertheless, according to eq. (10), this also implies much longer neutrino diffusion timescale and even those generated in the outermost layers of the core are trapped leading to the conversion of most of the envelope with little or no neutron ejection by neutrinos.

4.1. Ejection of neutrons

In the event that a significant fraction of the neutrinos manage to diffuse out of SQM before the conversion of the outermost layers of the hadronic envelope (e.g., for ρ_{uds} or core size such that $\tau < t_{coll}^{env}$ and small E_ν), they will interact with the electrons in the surface layer of quark matter and then with infalling neutrons. The $\nu-$(free) neutron absorption cross-section is larger by a factor 4 or more than elastic scattering at neutrino energies of 1 MeV, so one expects conversion of free neutrons to $p + e^-$. This cross-section is quite large (about $10^{-39}\mathrm{cm}^2$, implying a mean free path of a few hundred meters at envelope densities of 10^{13}g/cc). These sources of opacity (neutrino scattering off neutrons and electrons as well as absorption by nucleons) imply that for 1 MeV neutrinos the mean free path is roughly

$\lambda_{\rm env} < 5\,$km. Thus, it is possible that the neutrino energy is dumped in lower density nuclear matter surrounding the quark core, which we now proceed to estimate.

The total energy deposited by neutrinos in the nuclear envelope is given by

$$E_T = \tau \int_{R_1}^{R_2} Q_\nu(r) dr \tag{12}$$

where τ is a typical time for free-fall collapse, given approximately by $\tau = (R_2 - R_1)/v_c$ and v_c is the propagation speed of the detonation wave (we may take it to be the speed of sound in hot nuclear matter), R_1 and R_2 are the inner and outer radii of the nuclear envelope, and $Q_\nu(r)$ is the energy deposition rate per unit length at radius r. The energy loss is principally due to inelastic scattering of neutrinos on nucleons, with a typical mean free path of $\lambda_{\rm env} \simeq 1$ km. As we are interested in an order of magnitude estimate, we will approximate the radial dependence of the emissivity as tantamount to absorption by a thick medium of free scatterers, and indicate later how this approximation may be improved. For the simple exponential profile of energy loss, the emissivity from each successive volume element with increasing r drops as $\epsilon_\nu(r) = \epsilon_\nu(R_1){\rm e}^{-(r-R_1)/\lambda_{\rm env}}$. It follows that $Q_\nu(r) = 4\pi\, r^2\epsilon_\nu(R_1)(1 - {\rm e}^{-(r-R_1)/\lambda_{\rm env}})$. $\epsilon_\nu(R_1)$ is the neutrino energy impinging per unit volume per unit time at $r = R_1$, and is approximately equal to $L_\nu/(\frac{4}{3}\pi R_1^3)$, neglecting for the moment energy loss in the thin crust above quark matter. For the fiducial values $R_1 = 2$ km ($\tau = 1$ ms), $R_2 = 10$ km, $T_\nu \sim T_{\rm eff} = 10$ MeV ($\lambda_{\rm env} = 5$ m), we find

$$E_T = 1.5 \times 10^{51} {\rm ergs} \tag{13}$$

Energy deposition effectively stops when the infalling nuclear envelope is as thin as the typical mean free path. Our numbers can be improved with a more detailed treatment of the radial dependence of the energy deposition rate, and blast wave hydrodynamics. Although it appears that we have neglected neutrino emission processes from neutron matter itself, like the modified urca or neutrino bremsstrahlung, their inclusion does not change the order of magnitude estimate since for $T \leq 10$ MeV, the emissivity is comparable to or less than the emissivity from direct urca in quark matter. (This is because the T^8 dependence of these processes in neutron matter is more than compensated by smaller numerical pre-factors as compared to the direct urca in quark matter). We have also ignored nucleon correlations in dense matter and Pauli blocking effects, which effectively decrease the opacity, and would serve to soften the exponential fall-off of the emissivity. This implies that more energy could be deposited in the outer layers, resulting in somewhat increased mass ejection rates than we estimate below.

Mass ejection occurs when energy deposited (heating) by neutrinos in a shell of thickness

dr at radius r exceeds the gravitational energy density[3] induced by the mass $M(r) = M_1 + M_{\rm env}(r)$ with $M_1 = 4\pi/3\ \rho_{\rm uds} R_1^3$ as the core mass and $M_{\rm env}$ the envelope mass contained within r,

$$f(r) = \frac{\epsilon_\nu(R_1)(1 - {\rm e}^{-(r-R_1)/\lambda_{\rm env}})\ \tau}{\frac{G\ M(r)}{r}\rho_{\rm env}(r)} > 1\ . \tag{14}$$

Figure 4 shows the ratio ($f(x)$) between these two components versus radius ($x = r/R_1$) for different core sizes (R_1). Figure 5 shows the total mass ejected versus core radius R_1 for different core temperature, $T_{\rm eff}$. While ejection of inner layers (satisfying the $f(r) > 1$ condition) might be prevented by the overlaying envelope matter, for the cases we studied we find $f(r > r_{\rm min}) > 1$; $r_{\rm min}$ corresponds to $f(r = r_{\rm min}) = 1$. Thus the total ejected mass can simply be expressed as $M_{\rm ejec} = \int_{r_{\rm min}}^{R_2} 4\pi r^2 \rho_{\rm env}(r) dr$. We used our fiducial values, $R_2 = 10$ km, $T_\nu = T_{\rm eff}$ and $\rho_{\rm env}/\rho_{\rm uds} = 1/2$. We parametrize the density variation in the neutron rich envelope as $\rho_{\rm env.} \propto r^{-\alpha}$; we consider the $\alpha = 1$ and $\alpha = 2$ cases which reasonably describes the density variation over a radius range from 2-10 km, gathered from studies of neutron skin thickness in Pb nuclei (Horowitz & Piekarewicz 2001), as well as prior studies of the neutron matter equation of state (Strobel et al. 1997).

As can be seen in Figure 4 and Figure 5 the $\alpha = 1$ case requires extreme conditions (very small cores and high effective temperatures) for heating to ablate material. These extreme conditions also imply extreme mass ejection since the $\rho_{\rm env} \propto r^{-1}$ profile provides enough material in the outer layers. Specifically, the $\alpha = 1$ case requires temperatures above the 20 MeV range and core radius $R_1 < 1.2$ km for ejection to occur. Whether such small cores can provide the extreme temperatures following the collapse is questionable. As for the $\alpha = 2$ envelopes, mass ejection is triggered for core temperatures as low as 10 MeV. For example, for a core temperature $T_{\rm eff} \leq 15$ MeV, a 1.5 km core can ablate up to $0.01 M_\odot$ of envelope material.

We note that for the same envelope mass, $\alpha = 1$ envelopes are less dense than the $\alpha = 2$ ones. The neutrinos deposit less energy in such envelopes (larger $\lambda_{\rm env}$) explaining the need for extreme temperatures for ablation to be triggered. In both cases, however, we find that for small cores the neutrino flux is too small to account for any mass ejection while larger cores trap neutrinos long enough for the entire envelope to be converted to *(u,d,s)* matter.

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[3]The material is heated primarily via the charged-current absorption processes on free nucleon as $\nu_e n \rightarrow pe^-$ and $\bar{\nu_e} p \rightarrow ne^+$. A particle is unbound/ablated if the sum of its energies is positive; here we omitted the internal energy since previous studies have shown that it has practically no effect on the total amount of ejected material (Rosswog et al. 1999). In our simplified model this implies that most of the heat is converted to kinetic energy. In contrast to our estimate on the energy deposition by neutrinos, this introduces an overestimation of mass ejection, so that these effects compete in opposite directions.

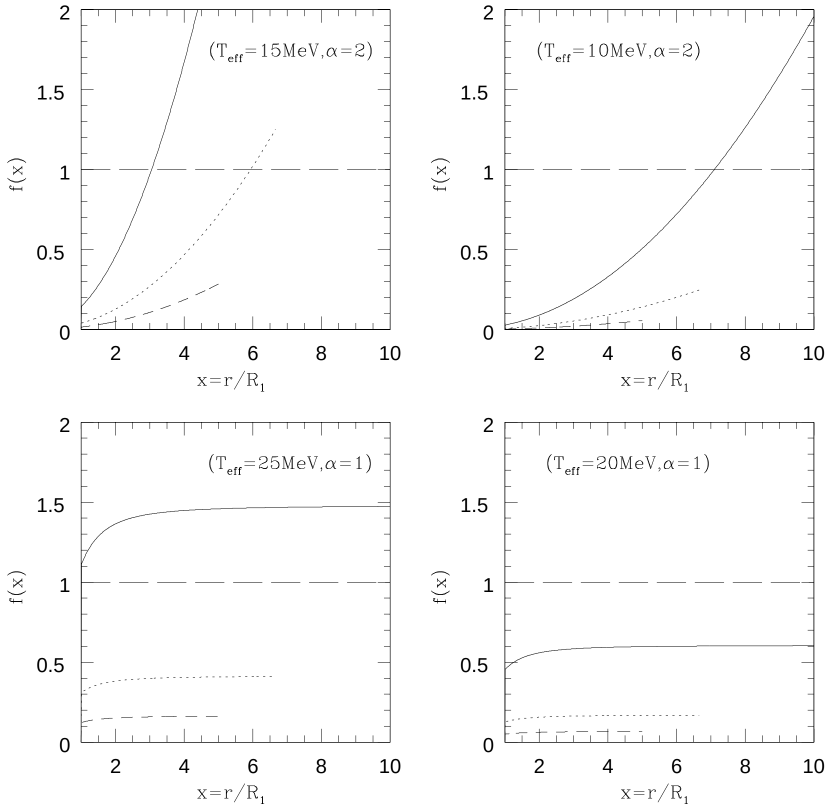

Fig. 4.— The ratio of neutrino energy density deposited to gravitational energy density versus radius of the envelope ($x = r/R_1$) for different core radius R_1. Mass ejection occurs when $f > 1$. Two different envelope density distribution $\rho_{\rm env} \propto r^{-\alpha}$ are shown, $\alpha = 2$ (upper panels) and $\alpha = 1$ (lower panels).

Finally, since the $\rho_{\rm env} \propto r^{-2}$ dependency is reflective of the outer radii of the neutron-rich envelope this favors the scenario where only the outermost layers will be ejected in QNe. Whether this ejected mass succeeds in turning into a wind and escape to infinity, or is later prevented by the fallback material remains to be confirmed.

4.2. Ejection of crust material ?

The crust material ($\rho_{\rm crust} < 10^{11}\,{\rm g\,cm^{-3}}$) subject to the Coulomb gap would remain suspended above the *(u,d,s)* and would most certainly be subject to the delayed neutrino burst. The amount of the crust ejected depends on the energy deposition rate which is

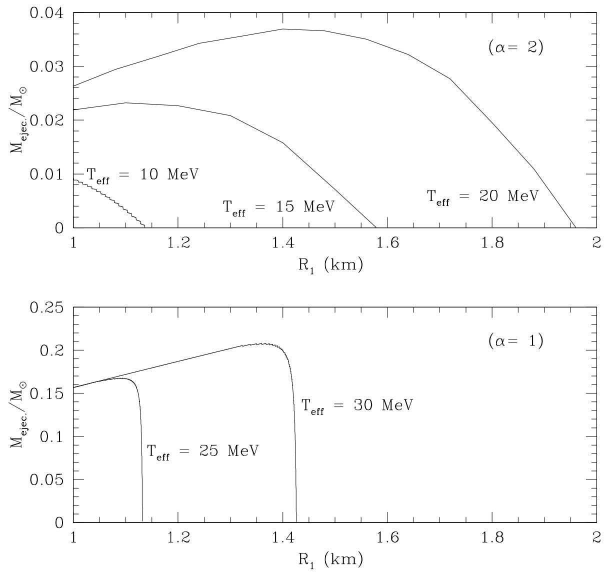

Fig. 5.— The total mass ejected (in solar units) versus the core size (R_1) for different core temperatures ($T_{\rm eff}$). Two different envelope density distribution $\rho_{\rm env} \propto r^{-\alpha}$ are shown, $\alpha = 2$ (upper panel) and $\alpha = 1$ (lower panel).

determined by the crust opacities; neutrino scattering off nucleons, nuclei and electrons as well as absorption by nucleons and nuclei all contribute, and the total opacity in the crust averaged for all neutrinos and anti-neutrinos is (Lamb & Pethick 1976)

$$\kappa_{\rm t} \simeq 1.2 \times 10^{-6} \left(\frac{\rho}{10^{10}\,{\rm g\ cm^{-3}}} \right) \left(\frac{kT_\nu}{10\,{\rm MeV}} \right)^2 {\rm cm}^{-1}, \tag{15}$$

implying a mean free path $\lambda = 1/\kappa_{\rm t}$ that is several orders of magnitude larger than the crust thickness. It is thus very unlikely that the crust will be blown off by neutrinos unless photons are at play as might be the case if color superconductivity sets in the core, (Ouyed & Sannino, (2002)). This also justifies neglecting the crust for the estimate of energy deposition in nuclear matter in the previous section.

4.3. Effects of color superconductivity

Due to the large critical temperatures associated to color superconductivity, the core may enter the color-flavor-locked (CFL) phase, with the surface in the 2-flavor superconducting phase (2SC). In that case, neutrino emission is dominated by the decay and scattering

of collective excitations, rather than free quarks. The neutrino opacity in the temperature range 10-30 MeV is a few meters, determined by neutrino absorption on the massless mode that describes baryon superfluidity (Reddy et al (2003)). For $T \sim 10$ MeV (still much smaller than the gap), the emission rates in CFL-paired matter are roughly three orders of magnitude smaller than unpaired quark matter, although the opacities are similar due to efficient scattering off Goldstone modes. Therefore, the neutrino flux from CFL matter is only about 1% of unpaired quark matter. Since the envelope is composed only of normal matter, a similar suppression is expected in the total mass ejected. Although the emission rate from the 2SC phase is only slightly reduced compared to normal quark matter, (beta decay of free light quarks of one color dominates), the fact that it occurs in a surface layer implies that the integrated luminosity receives contributions predominantly from CFL matter. The absence of Goldstone modes in the 2SC phase implies reduced opacities, so that the neutrinos from CFL quark matter can escape once they reach the quark matter surface. For temperatures of the order of the superconducting gap Δ, quasi-free excitations (Pair-breaking) can become dominant sources of neutrino emission and scattering (Jaikumar, Prakash, & Schaefer 2002; Kundu & Reddy 2004). The dominant neutrino rates then approach those in normal matter. Therefore, we conclude that the neutrino flux from the quark core, and hence the mass ejection is suppressed by at most 99% for $T \ll \Delta$ (effective quenching), and exponentially approaches the rate computed with normal quark matter as $T \rightarrow T_c$ (since the suppression of direct urca goes like exp(-Δ/T)). In light of the uncertainty in the exact densities at which different superconducting phases appear, and the value of the gaps themselves, these estimates are reasonable. The general conclusion is that the transition to a superconducting state will diminish the amount of energy deposition and mass ejection in the envelope through neutrinos.

4.4. Neutrino oscillations and matter effects

Recent experimental neutrino data provides evidence for neutrino oscillations (for a recent review, see e.g. Maltoni et al., 2004). Inside the dense core of the hybrid star, electron neutrinos (and antineutrinos) created in processes described by eqs. (2) are created as weak interaction eigenstates, but they propagate as the eigenstates of Hamiltonian, providing the mechanism for neutrino oscillations. Matter affects neutrino oscillations (Wolfenstein 1978, Mikheyev & Smirnov 1985) and may induce conversions between the neutrino states. These conversions occur at the so-called resonance densities (see e.g. Dighe & Smirnov, 2000). In the standard picture, there are two such resonance regions with densities of roughly $1000-10000$ g/cm^3 and $10-30$ g/cm^3. These densities are much lower than in the hadronic envelope, and therefore neutrino oscillations in matter do not influence our estimates of mass

ejection. Electron neutrinos (and anti-neutrinos) propagate within the quark star effectively in the same state that they were created.

In non-standard neutrino theories one may expect different effects. E.g. sterile neutrinos could, if created in conversion processes, carry away a significant amount of explosion energy and influence the dynamics of the QN. This kind of neutrino models have been extensively studied in connection with supernovae, and it is found that sterile neutrinos with masses in the keV-range may lead to conversions, but very heavy or light sterile neutrinos have no resonances inside the core or envelope (see e.g. Abazajian, Fuller & Patel (2001) and Keränen, Maalampi, Myyryläinen & Riittinen (2004)). We do not consider the possibility of sterile neutrinos further in this work.

To conclude, it appears that neutrino oscillations, masses and so-called matter effects do not change our results in the new standard neutrino physics picture.

5. Conclusion

We have studied the role of neutrinos in a quark nova, and found that they could deposit sufficient energy in the outer layers of the star to cause significant mass ejection of the nuclear envelope. This conclusion assumes a scenario in which the size of the quark core ($R_{\rm core} \leq 2$ km), is such that neutrinos can diffuse out and lose their energy in the nuclear envelope well before (u, d, s) conversion followed by shrinking of the star is completed. The mass ejection fraction of the outermost layers of the envelope is estimated using the energy deposition rate by neutrino-nucleon inelastic scattering. We predict on average up to $\sim 10^{-2} M_{\odot}$ of neutron rich material to be ejected during the explosion. This is suggestive of some interesting astrophysical implications such as r-process products injected into the inter-stellar medium (e.g., Freiburghaus et al. 1999 and references therein) and neutron-rich disk forming around newly born quark stars (e.g. Keränen & Ouyed 2003).

On the other hand, for the case when ($R_1 >> 1$ km) the conversion of the core into strange matter and the shrinking timescale into a dense quark object are faster than the neutrino diffusion time scale. The entire neutron star converts into a *(u,d,s)* object with no mass ejection.

Although we have mentioned two distinct physical scenarios, one where the envelope-core boundary is continuous, and another where they are separated by a macroscopic distance, the estimates and conclusions on neutrino transport and mass ejection are unaffected by this difference. However, they must be distinguished since the latter involves the interesting possibility of a low density region developing inside the star!

Clearly, to understand the complex energetics and dynamics involved in the QN explosion and the consequences on the surrounding environment, one needs the help of advanced numerical simulations where general relativistic effects can also be taken into account. Simplifying assumptions such as the stasis of the envelope while ejecting mass can be relaxed. The radial dependence of the energy deposition by neutrinos can be computed more accurately using Boltzmann equations for neutrino transport. In this work, we have outlined the basic physical picture of the quark nova and shown that mass ejection is feasible, based on dominant neutrino emissivities and opacities. Our future investigations are directed towards making this statement quantitatively precise by including the refinements of blastwave hydrodynamics and neutrino transport.

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